CEBAF

The Continuous Electron Beam Accelerator Facility

Theory Group Preprint Series

Additional copies are available from the authors.

The Southeastern Universities Research Association (SURA) operates the Continuous Electron Beam Accelerator Facility for the United States Department of Energy under contract DE-AC05-84ER40150

WM-93-106
CEBAF-TH-93-6

Isoscalar meson exchange currents and the deuteron form factors

Hiroshi Ito
Center for Nuclear Studies, Department of Physics
George Washington University
Washington, D.C. 20052

Franz Gross
Physics Department, College of William and Mary, Williamsburg, Virginia 23185
and
Continuous Electron Beam Accelerator Facility
12000 Jefferson Avenue, Newport News, Virginia 23606

The electromagnetic form factors for the $\rho\pi\gamma$ and $\omega\sigma\gamma$ vertices are calculated from quark loop diagrams which take the $q\bar{q}$ structure of the π, σ, ρ, and ω mesons into account. The resulting form factors decrease with increasing Q^2 (the square of the four-momentum of the off-shell photon) considerably more rapidly than the monopole form factors obtained from vector meson dominance. The implications of this behavior, which has a significant effect on the elastic electromagnetic form factors of deuteron, is discussed.

Typeset Using *REVTEX*

A number of experimental measurements, the most famous being the electrodisintegration of the deuteron at threshold, have established the existance of *isovector* meson exchange currents [1]. These currents are large because of the small mass and narrow width of the pion, and because the $\pi\pi\gamma$ coupling (equal to the electric charge) is comparable to the coupling of the photon to the nucleon itself. In contrast, the nature and size of *isoscalar* exchange currents is still an issue of some controversy. The $\rho\pi\gamma$ interaction certainly leads to an isoscalar exchange current, but because of the large mass and width of the ρ and the comparatively small size of the $\rho\pi\gamma$ coupling (to be discussed below), such a current is hard to distinguish from other short range interaction currents, including those which might arise from quark exchange forces [2]. The connection of the $\rho\pi\gamma$ coupling to the $A\dot{V}V$ anomaly does give this current a special status, however [3]. The introduction of isoscalar exchange currents derived from other photon couplings, such as the $\omega\sigma\gamma$ interaction, is more dubious. The very short range nature of this current, together with the phenomenological status of the σ, makes it hard to justify singling it out for special consideration. In this letter we discuss the evidence for the existance of such currents.

The simplest system in which to look for isoscalar exchange currents is the deuteron. The deuteron form factors have been calculated using a variety of relativistic schemes for treating the nuclear dynamics [4–8], and good deuteron wave functions can be derived from realistic models of the NN interaction based on meson field theory [9, 10]. The relativistic theory for the deuteron form factor is therefore fairly reliable. In the context of a Bethe-Salpeter one boson exchange (OBE) model of the nuclear force, the form factors can be calculated from only two contributions: the relativistic impulse approximation (RIA) in which the photon couples directly to one of the bound nucleons (shown in Fig. 1a), and the meson exchange current (MEC) contribution in which the photon couples to the exchanged mesons (shown in Fig. 1b). Because the deuteron is an isospin zero target, only the isoscalar MEC can contribute and, in the context of the OBE model, the $\rho\pi\gamma$ and $\omega\sigma\gamma$ currents are two likely candidates.

The $\rho\pi\gamma$ and $\omega\sigma\gamma$ exchange currents make small contributions to the magnetic and quadruple moments (which are the values of the magnetic and quadrupole form factors at $Q^2 = 0$, where Q^2 is the square of the four momentum tansferred by the photon), but these contributions are much less than $\frac{1}{2}$ % [14]. They are therefore significantly smaller than the relativistic corrections, which are about 5% for the magnetic moment and 1.5% for the quadrupole moment [9]. Thus the values of the static moments cannot be used to establish the existence of isoscalar exchange currents. However, as recently suggested by

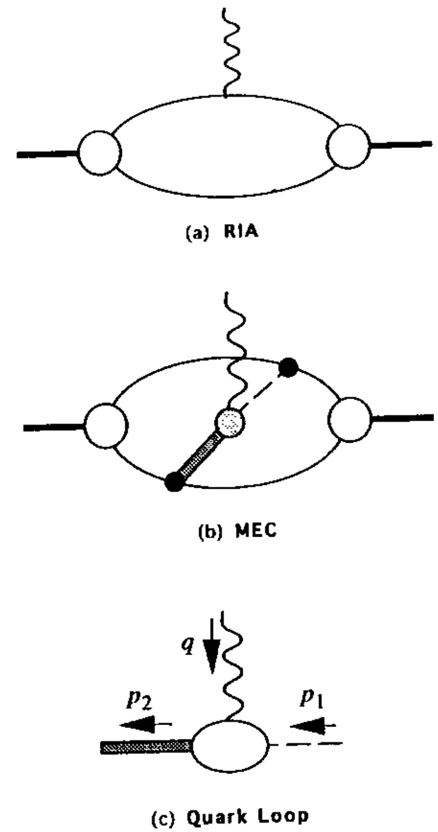

Figure 1: (a) The relativistic impulse approximation. (b) The $\rho\pi\gamma$ (or $\omega\sigma\gamma$) meson exchange current contribution. The ρ (or ω) is denoted by the wide shaded line and the π (or σ by the dashed line. (c) The quark loop contribution to the $\rho\pi\gamma$ (or $\omega\sigma\gamma$) vertex, represented by the lightly shaded circle in (b).

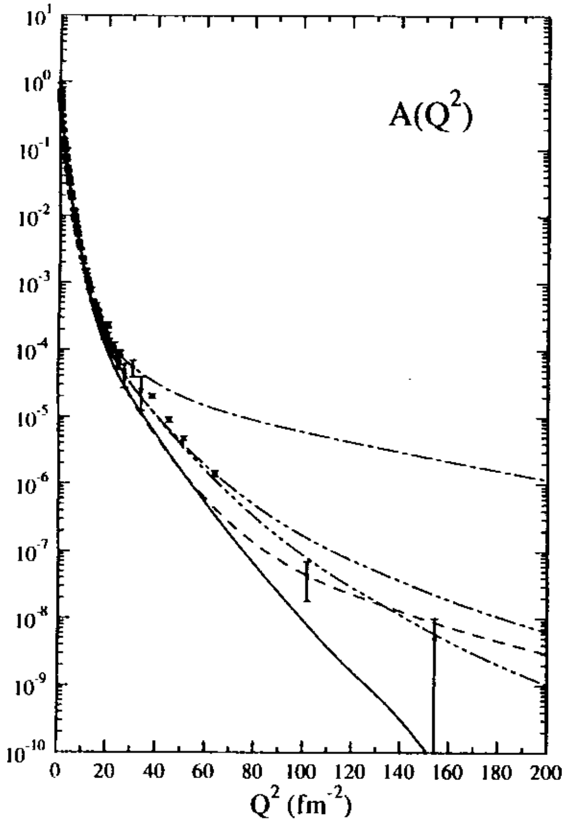

Figure 2a: Calculations of the A structure function under various assumptions. The solid line is the RIA of Hummel and Tjon, and the lines with *one, two, or three* short dashes are the RIA plus $\rho\pi\gamma$ exchange current with *no form factor*, the vector dominance form factor, or the quark loop form factor at the $\rho\pi\gamma$ vertex, respectively. (Addition of the $\omega\sigma\gamma$ exchange current with a quark loop form factor has a negligible effect.) The dashed line is the full calculation of Hummel and Tjon, which includs the RIA and $\rho\pi\gamma$ and $\omega\sigma\gamma$ exchange currents with vector dominance form factors.

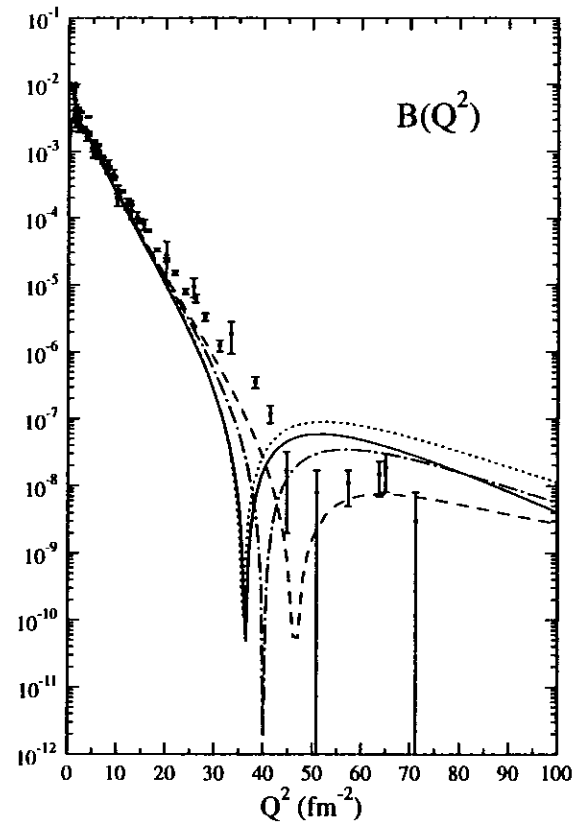

Figure 2b: Calculations of the B structure function under various assumptions. The solid line is the RIA of Hummel and Tjon, the dotted line is the RIA plus the $\rho\pi\gamma$ exchange current (with a quark loop form factor) and the dot-dashed line is the RIA plus *both* the $\rho\pi\gamma$ and $\omega\sigma\gamma$ exchange currents calculated with quark loop form factors. The dashed line is the full calculation of Hummel and Tjon (with VMD form factors).

Hummel and Tjon [14], the large Q^2 dependence of these form factors may be sensitive to these exchange currents, and it is this sensitivity which will be examined in this letter.

The differential cross section for the scattering of unpolarized electrons from deuterium can be written

$$\frac{d\sigma}{d\Omega} = \sigma_{\text{Mott}} \left(\frac{E'}{E}\right) \left[A(Q^2) + B(Q^2)\tan^2\left(\frac{\theta}{2}\right)\right] \tag{1}$$

where σ_{Mott} is the Mott cross section, E and E' are the incident and final electron energies, θ is the electron scattering angle, and $A(Q^2)$ and $B(Q^2)$ are structure functions related to the squares of the deuteron form factors. At high momentum transfers previous calculations of the RIA have underestimated $A(Q^2)$ by an order of magnitude at $Q^2 = 4\ GeV^2/c^2$ and also have failed to predict the correct location for the dip in $B(Q^2)$, as shown in Fig. 2. Because the MEC contributions provide a mechanism for sharing the incoming photon momentum equally between the two nucleons, they are expected to dominate at large momentum transfers [11–13], and recently Hummel and Tjon [14] used isoscalar meson exchange currents derived from the $\rho\pi\gamma$ and $\omega\sigma\gamma$ vertices to resolve the discrepancy at high Q^2. However, the size of these contributions depends critically on the Q^2 dependence of the form factors associated with the $\rho\pi\gamma$ and $\omega\sigma\gamma$ vertices; if no form factors are used (for example) the result from the $\rho\pi\gamma$ exchange current alone would overestimate $A(Q^2)$ by two order of magnitudes (Fig. 2a). In order to predict the high Q^2 dependence of the deuteron structure functions we must have a reliable estimate of the Q^2 dependence of the form factors at the $\rho\pi\gamma$ and $\omega\sigma\gamma$ vertices, and such an estimate should take into account the composite nature of the π, σ, ρ, and ω mesons. In this letter we calculate these form factors from quark loop diagrams which include the $q\overline{q}$ composite structure of the π, σ, ρ, and ω mesons, and discuss the implications of our results for the MEC contributions to the deuteron form factors.

In their calculation, Hummel and Tjon used a simple monopole form factor $f_{\rho\pi\gamma}(Q^2) = m_\omega^2/(m_\omega^2 + Q^2)$ where m_ω is the mass of the ω meson. Such a form factor is justified by the vector meson dominance (VMD) hypothesis. The applicability of VMD is a controversial topic in particle physics and using it to estimate the $\rho\pi\gamma$ form factor is particularly questionable if the photon momentum is large and space-like. Though it may work well at low Q^2, it is expected to deviate at high Q^2. Using asymptotic power counting based on perturbative QCD, Chernyak and Zhitnitsky [15] have shown that the $f_{\rho\pi\gamma}(Q^2) \sim Q^{-4}$, where the extra power in the fall off (the typical asymptotic meson form factor is $\sim Q^{-2}$) is due to the helicity flip of a quark. Apparently, at some momentum scale, the form factor must start to deviate from the monopole function. Our estimates, based on a quark loop calculation which takes the structure of the mesons in account, give a similar result.

To calculate the form factor at the $\rho\pi\gamma$ vertex we use a relativistic quark model of the pion and the rho which was previously used to calculate a variety of pion observables [16]. In this model the Bethe-Salpeter vertex function for the pion is taken to be $\Gamma_\pi(k,p) = \mathcal{N}_\pi\gamma^5\chi_f\chi_c D_S^{-1}(k^2)$, where k and p are the relative and total four momenta of $q\overline{q}$ pair which couple to the pion, $D_S(k^2) = k^2 - \Lambda^2$ with Λ an adjustable parameter, $\chi_c = I/\sqrt{N_c}$ the color wave function, and $\chi_f = \tau^+$ the flavor wave function of the π^+. The function $D_S^{-1}(k^2)$ is a parameterization of the momentum distribution of the $q\overline{q}$ pair in the pion. The normalization constant, $\mathcal{N}_\pi$, is fixed by the requirement that the charge form factor of pion, $F_\pi(Q^2)$, be normalized to $F_\pi(0) = 1$. In Ref. [16] the form factor and low-energy observables of the pion were very well reproduced by choosing the quark mass $m_q = 248$ MeV and $\Lambda = 450$ MeV. This model also successfully described the Q^2 dependence of both the pion form factor and the form factor in $\gamma^* + \pi^0 \to \gamma$ process [17], which has been recently measured at space-like virtual photon (γ^*) momenta in the range $Q^2 = 1 \sim 3$ GeV$^2/c^2$ [18]. For the ρ meson vertex we choose the form $\Gamma_\rho^\nu(k,p) = \mathcal{N}_\rho G^\nu(p)\chi_f\chi_c D_V^{-1}(k^2)$, where $G^\nu(p) = [\gamma^\nu - \not{p}p^\nu/p^2]$, and the normalization constant, $\mathcal{N}_\rho$, is calculated from the residue of the $q\overline{q}$ scattering amplitude in the vector channel [16]. A dipole momentum distribution $D_V(k^2) = (k^2 - \Lambda_1^2)(k^2 - \Lambda_2^2) \sim k^4$ was chosen in order to insure that the integrals involving the rho converge. Using the same quark mass, the choice $\Lambda_1 \sim 600$ MeV and $\Lambda_2 \sim 1000$ MeV fits the empirical values of the $\gamma\rho$ coupling and the rho width. In this calculation, in order to avoid threshold singularities, we found it convenient to use a larger quark mass of $m_q \simeq 390$ MeV $> m_\rho/2$. Using this larger quark mass changes the fitted observables by only about 25%, which is sufficiently close for our estimates.

Now we use this model to calculate the $\rho\pi\gamma$ vertex, which is defined by [12],

$$\langle\rho(p_2)|J^\mu|\pi(p_1)\rangle = \mathcal{G}(Q^2)\epsilon^{\mu\nu\alpha\beta}p_{1\alpha}q_\beta\epsilon_\nu \tag{2}$$

where the antisymmetric tensor $\epsilon^{\mu\nu\alpha\beta}$ assures electromagnetic gauge invariance, q ($q^2 = -Q^2$) is the four momentum of the virtual photon, and ϵ_ν is the polarization vector of the ρ meson. The dependence on the momentum of photon is expressed by $\mathcal{G}(Q^2) = -i\frac{e}{m_\rho}g_{\rho\pi\gamma}f_{\rho\pi\gamma}(Q^2)$, where the form factor $f_{\rho\pi\gamma}(Q^2)$ is normalized to unity at $Q^2 = 0$ and $g_{\rho\pi\gamma}$ is the coupling constant. We calculated the product of this form factor and the coupling constant from the quark loop diagram shown in Fig1c, which gives

$$\langle\rho(p_2)|J^\mu|\pi(p_1)\rangle = e_q \int \frac{d^4k}{(2\pi)^4} Tr\Big\{\epsilon_\nu \overline{\Gamma}^\nu_\rho(K_f;p_2) S\left(k+\frac{q}{2}\right)\ \gamma^\mu\, S\left(k-\frac{q}{2}\right)$$
$$\times \Gamma_\pi(K_i;p_1) S\left(k-\frac{p_1+p_2}{2}\right)\Big\} + \left(e_{\overline{q}}\ \text{term}\right) \quad (3)$$

where $K_f = k - p_1/2$, $K_i = k - p_2/2$, $\Gamma_\pi(k;p)$ and $\Gamma^\nu_\rho(k;p)$ are the vertex functions of π and ρ mesons, $S(k)$ is the quark propagator, and e_q is the quark charge operator which takes on values of $e_u = \frac{2}{3}$ and $e_d = -\frac{1}{3}$. The same expression was used for the amplitude for the $\gamma^* + \pi^0 \to \gamma$ process except that in that calculation $\overline{\Gamma}^\nu_\rho(k;p)$ was replaced by the photon-quark vertex, $e_q\gamma^\nu$. The measured form factor for the $\gamma^* + \pi^0 \to \gamma$ exhibits a monopole behavior of the form $f_{\gamma^*+\pi^0\to\gamma}(Q^2) = \lambda^2/(\lambda^2+Q^2)$, with $\lambda = 748 \pm 30$ MeV. Because the $\rho\pi\gamma$ vertex substitutes the vertex function of the ρ in place of the point-like quark-photon coupling, the fall off of the corresponding $f_{\rho\pi\gamma}(Q^2)$ should be more rapid than a monopole function. Our calculated form factor, shown in Fig. 3, falls off much more rapidly at large Q^2 than a monopole; it is about a factor of 3 times smaller at $Q^2 = 200$ fm^{-2}. The calculated value of $g_{\rho\pi\gamma}$ is 0.71, in fair agreement with the recently measured value of 0.56 [19].

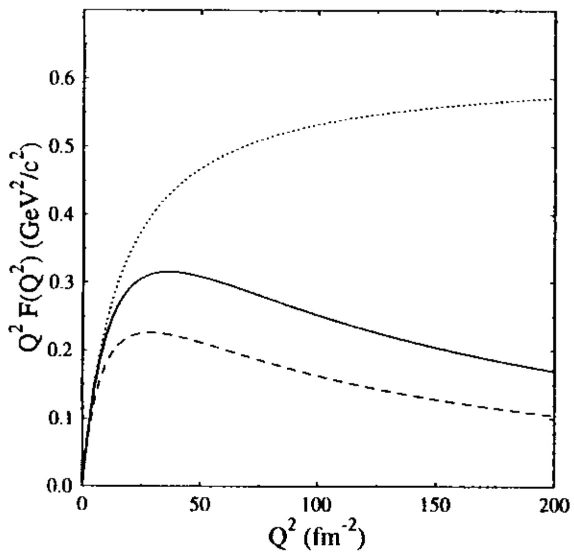

Figure 3: The quark loop $\rho\pi\gamma$ (solid line) and $\omega\sigma\gamma$ (dashed line) form factors multiplied by Q^2. The vector dominance form factor (dotted line) is shown for comparison.

We now turn to a discussion of the $\omega\sigma\gamma$ exchange current. Here we are less able to present a reliable calculation, partly because it is unclear whether to take the "σ" to be the observed $I = 0$ resonance with a mass in the vicinity of 1000 MeV, or the phenomenological two pion enhancement which plays a role in all OBE models of the nuclear force and which has a mass in the vicinity of 500 MeV. In our estimate we will follow Ref[14] and assume the latter. We will also assume this sigma to be the chiral partner of the pion [20, 21], which suggests, in the spirit of the NJL model [21], that the $\sigma q\overline{q}$ vertex can be obtained from the $\pi q\overline{q}$ vertex by replacing the factors of $\gamma^5\chi_f$ by the unit matrix (so that the momentum dependence is the same). Similarily, we will use the same parameterization for the $\omega q\overline{q}$ vertex as we did for the $\rho q\overline{q}$ vertex. With these assumptions, we will calculate the $\omega\sigma\gamma$ interaction from the quark loop diagram in Fig 1c.

This calculation is complicated by the fact that the simple loop diagram is, by itself, not gauge invariant. To obtain a gauge invariant result we must take account of the interaction currents induced by the nonlocality of the meson-$q\overline{q}$ vertices. Given models of $\omega q\overline{q}$ and $\sigma q\overline{q}$ vertices, $\Gamma_\omega(k,p)$ and $\Gamma_\sigma(k,p)$, the method of minimal substitution [22] gives explicit forms for these interaction currents,

$$\mathcal{J}^\mu_{\sigma q\overline{q}}(k;p,q) = e_u[\Gamma_\sigma(k-\frac{q}{2},p) - \Gamma_\sigma(k,p)]\frac{[4k-q]^\mu}{q\cdot[4k-q]} + \left(e_{\overline{d}}\ \text{term}\right), \quad (4a)$$

$$\overline{\mathcal{J}}^\mu_{\omega q\overline{q}}(k;p,q) = e_u\,\epsilon_\nu[\overline{\Gamma}^\nu_\omega(k,p) - \overline{\Gamma}^\nu_\omega(k+\frac{q}{2},p)]\frac{[4k+q]^\mu}{q\cdot[4k+q]} + \left(e_{\overline{d}}\ \text{term}\right), \quad (4b)$$

where the ($e_{\overline{d}}$ term) is obtained by replacing $q \to -q$ in the e_u term. Including these contributions the $\omega\sigma\gamma$ amplitude is given by

$$\langle\omega(p_2)|J^\mu|\sigma(p_1)\rangle = e_q \int \frac{d^4k}{(2\pi)^4} Tr\Big\{\epsilon_\nu \overline{\Gamma}^\nu_\omega(K_f;p_2) S\left(k+\frac{q}{2}\right)\ \gamma^\mu\, S\left(k-\frac{q}{2}\right)$$
$$\times \Gamma_\sigma(K_i;p_1) S\left(k-\frac{p_1+p_2}{2}\right)\Big\} + \left(e_{\overline{q}}\ \text{term}\right)$$
$$+ \int \frac{d^4k}{(2\pi)^4} Tr\left\{S\left(k-\frac{p_1+p_2}{2}\right)\ \epsilon_\nu \overline{\Gamma}^\nu_\omega(K_f;p_2) S\left(k+\frac{q}{2}\right) \mathcal{J}^\mu_{\sigma q\overline{q}}(K_f;p_1,q)\right\}$$

$$+\int \frac{d^4k}{(2\pi)^4} Tr \left\{ \overline{\mathcal{J}}^{\mu}_{\omega q\bar{q}}(K_i;p_2,q) S\left(k-\frac{q}{2}\right) \Gamma_\sigma(K_i;p_1) S\left(k-\frac{p_1+p_2}{2}\right)\right\} \quad (5)$$

and the gauge invariance of the amplitude, $\langle\omega^\nu(p_2)|q_\mu J^\mu|\sigma(p_1)\rangle = 0$, follows from the use of the Ward-Takahashi identity at the quark-photon vertices. Explicit evaluation of this expression gives two independent terms:

$$\begin{aligned}\langle\omega(p_2)|J^\mu|\sigma(p_1)\rangle &= \mathcal{G}^{(a)}(Q^2)\Big[(q\cdot p_1)\epsilon^\mu - (\epsilon\cdot q)p_1^\mu\Big] \\ &+ \mathcal{G}^{(b)}(Q^2)\Big[(q\cdot p_2)\epsilon^\mu - (\epsilon\cdot q)p_2^\mu\Big] \qquad (6)\end{aligned}$$

where $\mathcal{G}^{(a)}(Q^2)$ and $\mathcal{G}^{(b)}(Q^2)$ are independent functions of Q^2. These two terms also appear in models in which the sigma is treated as a composite two pion system [23].

Numerical evaluation of Eq. (5), shows that the two independent functions $\mathcal{G}^{(a)}(Q^2)$ and $\mathcal{G}^{(b)}(Q^2)$ have opposite sign and are nearly equal in magnitude (the sum of these two functions is less than 20% of their difference). In addition, the sum vanishes if $m_\omega \simeq m_\sigma$ and decreases more rapidly with Q^2 than the difference. For all of these reasons we approximate the $\omega\sigma\gamma$ vertex by

$$\langle\omega(p_2)|J^\mu|\sigma(p_1)\rangle = \frac{e}{m_\omega} g_{\omega\sigma\gamma} f_{\omega\sigma\gamma}(Q^2)\Big[q^2\epsilon^\mu - (\epsilon\cdot q)q^\mu\Big] \qquad (7)$$

where $\frac{e}{m_\omega} g_{\omega\sigma\gamma} f_{\omega\sigma\gamma}(Q^2) = \mathcal{G}^{(a)}(Q^2) = -\mathcal{G}^{(b)}(Q^2)$. The coupling constant is defined by the normalization $f_{\omega\sigma\gamma}(0) = 1$. We emphasize that Eq. (7) is only an approximation to the full result given in Eq. (6), and that its tensor structure, $q^2\epsilon^\mu - (\epsilon\cdot q)q^\mu$, differs from $(q\cdot p_1)\epsilon^\mu - (\epsilon\cdot q)p_1^\mu$, the result conventionally used. However the integral over the exchange current operator (Fig. 1b) is approximately symmetric under the replacement $p_1 \leftrightarrow -p_2$, and the two tensor forms are therefore approximately equivalent. Our numerical result for $f_{\omega\sigma\gamma}(Q^2)$ is shown in Fig. 3. Note that the factorization assumption, $f_{\omega\sigma\gamma}(Q^2) = constant \times f_{\rho\sigma\gamma}(Q^2)$, works approximately at large Q^2 ($Q^2 > 3$ GeV^2/c^2), as suggested by power counting. Factorization does not work at low Q^2, where the form factors are strongly affected by the difference in the matrix structure of the $\rho\pi\gamma$ and $\omega\sigma\gamma$ loops.

Fig.2 shows the effects of different models for $f_{\rho\pi\gamma}(Q^2)$ and $f_{\omega\sigma\gamma}(Q^2)$ on the deuteron structure functions $A(Q^2)$ and $B(Q^2)$. Fig. 2a shows that the large Q^2 dependence of $A(Q^2)$ is quite sensitive to $f_{\rho\pi\gamma}(Q^2)$; the differences of a factor ~ 10 at Q^2= 200 fm^{-2} are a consequence of differences in $f_{\rho\pi\gamma}(Q^2)$. At these momentum transfers $A(Q^2)$ depends entirely on the MEC contribution. Because Hummel and Tjon used a VMD form factor, their $\rho\pi\gamma$ exchange current was too

large at high Q^2, and they needed another exchange current to cancel it. The $\omega\sigma\gamma$ MEC has the opposite sign, and assuming factorization $[f_{\omega\sigma\gamma}(Q^2) = f_{\rho\pi\gamma}(Q^2)]$ and $g_{\rho\pi\gamma} = -g_{\omega\sigma\gamma} = 0.56$ results in a strong cancellation of these MECs, as shown in the figure (the curve shown here is a corrected version of the original curve published in Ref[14]). Using QL form factors, the $\omega\sigma\gamma$ MEC is smaller than the $\rho\pi\gamma$ MEC and the cancellation is much less sensitive. In this case the $\omega\sigma\gamma$ MEC is very small, but the $\rho\pi\gamma$ MEC current gives better agreement with the data.

Next, consider $B(Q^2)$ shown in Fig. 2b. The $\rho\pi\gamma$ MEC makes only a small contribution to $B(Q^2)$, and Hummel and Tjon found that they could obtain agreement with the data [24] only by introducing the $\omega\sigma\gamma$ MEC. Note that the position of the diffraction minimum is very sensitive to the choice of the $\omega\sigma\gamma$ form factor. While good agreement was obtained with the VMD model, the more realistic QL model does not succeed in fitting the data. We conclude that $B(Q^2)$ structure function is still unexplained.

In conclusion, we emphasize that the MEC contributions are extremely sensitive to the form factors at the photon-meson vertices. The predictions for the $A(Q^2)$ structure function for $Q^2 > 2$ GeV^2/c^2 depend strongly on these form factors. Furthermore, these predictions are also sensitive to the (unknown) size of the neutron charge form factor (see Ref[4]). Taking all of these considerations into account we find that the A structure function could easily be explained by some combination of a $\rho\pi\gamma$ exchange current and an enhanced neutron charge form factor, but that, in the absence of measurements of the neutron charge form factor, it is difficult to fix this combination. The situation is different for the B structure function, where the contributions from *both* the $\rho\pi\gamma$ exchange current and the neutron charge form factor are negligible. However, attempts to use an $\omega\sigma\gamma$ exchange current to explain the B form factor are not successful unless an unrealistically hard form factor is used. We conclude that the B form factor is not explained by present models.

ACKNOWLEDGMENTS

The authors are deeply indebted to Prof. John Tjon and Dr. E. Hummel for supplying their nuclear matrix elements of the $\rho\pi\gamma$ and $\omega\sigma\gamma$ curent operators. H.I acknowledges useful discussions with Prof. S. J. Wallace, Dr. Neal Devine and Prof. Mark Strickman. The work of H.I is partly supported by the Department of Energy Grant DE-FG05-86-ER40270 for the Center for Nuclear Studies, at the George Washington University. The work of F.G is partly supported by DOE Grant DE-FG05-88ER40435.

REFERENCES

[1] D. O. Riska and G. E. Brown, Phys. Lett. **38B**, 193 (1972). M. Bernheim, *et. al.*, Nucl. Phys. **A365**, 349 (1981).

[2] M. Oka and K. Yazaki, Phys. Lett. **90B**, 41 (1980); S. Takeuchi and K. Yazaki, Nucl. Phys. **A438**, 605 (1985); Y. Yamauchi and M. Wakamatsu, Nucl. Phys. **A457**, 621 (1986); H. Ito and A. Faessler, Nucl. Phys. **A470**, 626 (1987); T. de Forest and P. J. Mulders, Phys. Rev. D **35**, 2849 (1987); H. Ito and L. S. Kisslinger Phys. Rev. C **40**, 887 (1989); A. E. L. Dieperink and P. J. Mulders, Nucl. Phys. **A497**, 253c (1989).

[3] E. Nyman and D. O. Riska, Phys. Rev. Lett. **57**, 3007 (1986); Nucl. Phys. **A468**, 473 (1987). Wakamatsu and Weise, Nucl. Phys. **A477**, 559 (1988).

[4] R. G. Arnold, C. Carlson and F. Gross, Phys. Rev. C **21**, 1426 (1980).

[5] J. A. Tjon and M. J. Zuilhof, Phys. Lett. **84B**, 31 (1979); M. J. Zuilhof and J. A. Tjon, Phys. Rev. C **22**, 2369 (1980).

[6] R. Blankenbecler and R. Suger, Phys. Rev. **142** ,1051 (1966).

[7] P. L. Chung, F. Coester, B. D. Keister and W. N. Polyzou, Phys. Rev. C **37**, 2000 (1988).

[8] L. L. Frankfurt, I. L. Grach, L. A. Kondratyuk and M. I. Strickman, Phys. Rev. Lett. **62**, 387 (1989).

[9] F. Gross, J. W. Van Orden and K. Holinde, Phys. Rev. C **45**, 2094 (1992).

[10] R. Machleidt, K. Holinde and Ch. Elster, Phys. Rep. **149**, 1 (1987).

[11] R. J. Adler and S. D. Drell, Phys. Rev. Lett. **13**, 349 (1964).

[12] M. Chemtob, E. Moniz and M. Rho, Phys. Rev. C **10**, 344 (1974).

[13] M. Gari and H. Hyuga, Nucl. Phys. **A264**, 409 (1976).

[14] E. Hummel ans J. A. Tjon, Phys. Rev. Lett. **63**, 1788 (1989); Phys. Rev. C **42**, 423 (1990).

[15] V. L. Chenyak and A. R. Zhitnitsky, Phys. Rep. **112**, 173 (1984).

[16] H. Ito, W. Buck and F. Gross, Phys. Rev. C **45**, 1918 (1992); **43**, 2483 (1991).

[17] H. Ito, W. Buck and F. Gross, Phys. Lett. **287B**, 23 (1992).

[18] H. J. Behrend *et al.* Z. Phys. **C49**, 401 (1991).

[19] D. Berg *et al.*, Phys. Rev. Lett. **44**, 706 (1980).

[20] R. Delbourgo and M. D. Scadron, Phys. Rev. Lett. **48**, 379 (1982).

[21] Y. Nambu and G. Jona-Lasinio, Phys. Rev. **122**, 345 (1961).

[22] K. Ohta, Phys. Rev. C **40**, 1335 (1989).

[23] S. M. Berman and S. D. Drell, Phys. Rev. **133**, 791 (1964).

[24] R. G. Arnold *et al*, Phys. Rev. Lett. **58** , 1723 (1987).